

Cooling of Neutron Stars. Hadronic Model. [★]

D. Blaschke^{1,2}, H. Grigorian^{1,3}, and D.N. Voskresensky^{4,5}

¹ Fachbereich Physik, Universität Rostock, Universitätsplatz 1, D-18051 Rostock, Germany

email: david.blaschke@physik.uni-rostock.de

² Bogoliubov Laboratory for Theoretical Physics, Joint Institute for Nuclear Research, 141980 Dubna, Russia

³ Department of Physics, Yerevan State University, Alex Manoogian Str. 1, 375025 Yerevan, Armenia

email: hovik@darss.mpg.uni-rostock.de

⁴ Gesellschaft für Schwerionenforschung mbH, Planckstr. 1, D-64291 Darmstadt, Germany

⁵ Moscow Institute for Physics and Engineering, 115409 Moscow, Russia

email: d.voskresensky@gsi.de

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Abstract. We study the cooling of isolated neutron stars. The main cooling regulators are introduced: EoS, thermal transport, heat capacity, neutrino and photon emissivity, superfluid nucleon gaps. Neutrino emissivity includes main processes. A strong impact of medium effects on the cooling rates is demonstrated. With taking into account of medium effects in reaction rates and in nucleon superfluid gaps modern experimental data can be well explained.

Key words. dense baryon matter, neutron stars, medium effects, pion softening, nucleon gaps, heat transport

1. Introduction

The Einstein Observatory was the first that started the experimental study of surface temperatures of isolated neutron stars (NS). Upper limits for some sources have been found. Then ROSAT offered first detections of surface temperatures. Next *X-ray* data came from Chandra and XMM/Newton. Appropriate references to the modern data can be found in recent works by Tsuruta et al. 2002, Tsuruta 2004, Kaminker et al. 2001, Yakovlev et al. 2003a, devoted to the analysis of the new data. More upper limits and detections are expected from satellites planned to be sent in the nearest future. In general, the data can be separated in three groups. Some data show very “*slow cooling*” of objects, other demonstrate an “*intermediate cooling*” and some show very “*rapid cooling*”. Now we are at the position to carefully compare the data with existing cooling calculations.

2. Existing NS cooling scenarios

Let us briefly point out some important achievements on the way to the present understanding of the NS cooling problem. Theoretical study of the NS cooling has been started long ago in pioneering works of Tsuruta & Cameron 1965

and Bahcall & Wolf 1965. It has been argued that the one-nucleon so called direct Urca (DU) process, as $n \rightarrow pe\bar{\nu}$, is forbidden up to sufficiently high density and the main rôle plays the two-nucleon so called modified Urca (MU) process, like $nn \rightarrow npe\bar{\nu}$ and $np \rightarrow ppe\bar{\nu}$. As the result of many works the so called “*Standard scenario*” of cooling emerged. It includes the neutrino cooling stage $t \lesssim 10^5$ yr and the photon cooling era, $t \gtrsim 10^5$ yr. The MU process and the nucleon-nucleon bremsstrahlung (NB), as $nn \rightarrow nn\nu\bar{\nu}$ and $np \rightarrow np\nu\bar{\nu}$, were carefully recalculated by Friman & Maxwell 1979 using the free one-pion exchange. The emissivity of the most efficient channel of the MU process $nn \rightarrow npe\bar{\nu}$ is given by

$$\varepsilon_\nu[\text{MU}] \sim 10^{22} (m_{\text{MU}}^*(n))^4 \left[\frac{n_e}{n_0} \right]^{1/3} \xi_{nn} \xi_{pp} T_9^8 \frac{\text{erg}}{\text{cm}^3 \text{ sec}} \quad (1)$$

$T_9 = T/10^9 \text{K}$, n_e is the electron density and $n_0 \simeq 0.16 \text{ fm}^{-3}$ is the ordinary nuclear density at saturation, $(m_{\text{MU}}^*(n))^4 = (m_n^*/m_N)^3 (m_p^*/m_N)$, m_n^* and m_p^* are effective neutron and proton masses, m_N is the nucleon mass in vacuum, n is the nucleon density, $(m_{\text{MU}}^*(n))^4 (n_e/n_0)^{1/3} \sim 0.1$ for $n \sim n_0$.

After Migdal 1959 it became clear that NS in the late cooling stage for $T < T_{cn}$ and $T < T_{cp}$ are neutron and proton superfluids. The nucleon superfluidity was incorporated in the “*Standard scenario*” by Maxwell 1979, who used relevant combinations of the suppression factors

$$\xi_{ii} \simeq \exp(-\Delta_{ii}/T), \quad T < T_{ci}; \quad (2)$$

Send offprint requests to: D. Blaschke

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$\xi_{nn}\xi_{pp}$ for the emissivity of MU $nn \rightarrow npe\bar{\nu}$ and NB $np \rightarrow np\nu\bar{\nu}$ processes, ξ_{nn}^2 for NB $nn \rightarrow nn\nu\bar{\nu}$, ξ_{pp}^2 for MU $np \rightarrow ppe\bar{\nu}$, for $T < T_{cn}$ and $T < T_{cp}$, Δ_{ii} is the gap, $i = n$ or p .

Typically the NB emissivity is by an order of magnitude smaller than that for MU. The pre-factor $\sim 10^{20} \div 10^{21}$ and the temperature dependence T_9^8 are typical for the phase space volume of two-nucleon processes (without inclusion of medium effects, see below).

First cooling calculations were based on the assumption of an isothermal core. The relevance of the thermal conductivity for the first $10^2 \div 10^3$ yr was demonstrated by Nomoto & Tsuruta 1981. A discussion of the achievements of early works can be found in the review by Tsuruta 1979 and in the book by Shapiro & S.A. Teukolsky 1983.

The “*Standard scenario*” allows to explain the *slow cooling* but fails to cook the *intermediate cooling* and *fast cooling* of some neutron stars. To explain the latter different efficient direct Urca-like processes have been suggested: the proper DU process, reinvestigated by Lattimer et al. 1991 and the DU process going on hyperons (HDU), see the same work; pion condensation (PU) processes by Maxwell et al. 1977; kaon condensation (KU) processes by Brown et al. 1988, Tatsumi 1988, similar to PU; and quark DU processes (QDU) by Iwamoto 1982. All these processes show up for baryon densities larger than the corresponding critical ones $n > n_c^{\text{DU}}, n_c^{\text{PU}}, n_c^{\text{KU}}, n_c^{\text{HDU}}, n_c^{\text{QDU}} \sim 1.5 \div 6n_0$ according to different model dependent estimates. Roughly

$$\varepsilon_\nu[\text{DU}] \sim 10^{27} (m_{\text{DU}}^*)^2 \left[\frac{n_e}{n_0} \right]^{1/3} T_9^6 \times \min[\xi_{nn}, \xi_{pp}] \frac{\text{erg}}{\text{cm}^3 \text{ sec}}, \quad (3)$$

where $(m_{\text{DU}}^*)^2 = (m_n^* m_p^*)/m_N^2$. The factors $10^{26} \div 10^{29}$ and the behavior $\sim T_9^6$ are typical for all mentioned one-nucleon processes. Thus the “*Standard scenario+exotics*” has been considered as a scenario with minimum exotics: some stars have $n < n_c$ and cool slowly whereas some have $n > n_c$ in a part of their interiors and cool very fast.

However one type of the non-exotic processes was completely forgotten already in the “*Standard scenario*” without any justification for that. These are the so called neutron pair breaking and formation (nPBF) and proton pair breaking and formation (pPBF) processes, as they were named by Schaab et al. 1997. The emissivity of the nPBF process was first calculated by Flowers et al. 1976 for the case of the $1S_0$ neutron pairing. However, their asymptotic expression for the emissivity $\varepsilon[\text{nPBF}] \sim 10^{20} T_9^7 \exp(-2\Delta_{nn}/T)$ for $T \ll \Delta_{nn}$, as follows from expression (1b) of their work and from their rough asymptotic estimate of the integral (see below (13b)), shows neither the large one nucleon phase space factor ($\sim 10^{29}$) nor the appropriate temperature behaviour (note that the full analytic expression of this work is quite correct). The numerical underestimation of the rate by an order of magnitude and the very rough asymptotic expression used by Flowers et al. 1976 became the reason

that this important result was overlooked for many years. Voskresensky & Senatorov 1987 rediscovered the nPBF process $n \rightarrow n\nu\bar{\nu}$ and introduced the pPBF process $p \rightarrow p\nu\bar{\nu}$ with correct asymptotic behavior of the emissivity

$$\varepsilon_\nu[\text{iPBF}] \sim 10^{29} m_{\text{iPBF}}^* \left[\frac{p_{Fi}(n)}{p_{Fn}(n_0)} \right] \left[\frac{\Delta_{ii}}{\text{MeV}} \right]^7 \left[\frac{T}{\Delta_{ii}} \right]^{1/2} \times \xi_{ii}^2 \frac{\text{erg}}{\text{cm}^3 \text{ sec}}, \quad (4)$$

for $T \ll \Delta_{ii}$, $p_{Fi}(n)$ is the neutron/proton Fermi momentum, $i = n, p$, $m_{\text{iPBF}}^* = m_i^*/m_N$. Eq. (4) shows a huge one-nucleon phase space factor and a quite moderate T dependence of the pre-factor (note that the value of the pre-factor in case of $1S_0$ neutron pairing evaluated by Voskresensky & Senatorov 1987 contains a mistake that however yields only a non-essential correction factor of order of one to their final result; a larger uncertainty comes from not too well known medium modified vertices, see discussion by Voskresensky 2001 and corrected expressions there). Medium modifications of vertices (cf. Ward-Takahashi identity) were incorporated in all processes. This is very important point since medium modified vertices yield an enhancement factor up to 10^2 for pPBF compared with the result one would get using vacuum vertices. This factor $\sim 10^2$ arises since the process may occur through nn^{-1} and ee^{-1} correlations, -1 symbolizes the particle hole, with subsequent production of $\nu\bar{\nu}$ from the n and e rather than from a strongly suppressed vacuum vertex of $p\nu\bar{\nu}$, see Voskresensky & Senatorov 1987, Voskresensky et al. 1998, Leinson 2000, Voskresensky 2001. The efficiency of the PBF rates was related by Voskresensky & Senatorov 1987 to the value of the pairing gap and it was indicated that these processes can play an important rôle in the cooling scenario. Note that with taking into account medium effects the estimate (4) is roughly valid for all the cases for $1S_0$ neutron and proton pairings as well as for the $3P_2$ neutron pairing, except the specific case, when the projection of the pair momentum onto the quantization axis is $|m_J| = 2$. In the latter case, where the gap has zero's in some points of the Fermi surface, the emissivity would not be exponentially suppressed, cf. Voskresensky & Senatorov 1987. However conditions, when such a possibility might be realized, are not clear.

The inclusion of the PBF processes in the cooling code has first been accomplished by Schaab et al. 1997 (where also for the first time the reference to the pioneering work of Flowers et al. 1976 appeared) with the result that variation of gaps allows to develop a “*scenario for intermediate cooling*” covering the region of both “*slow and intermediate cooling*”. Some subsequent works rediscovered the pPBF process with vacuum vertices thus ignoring medium effects but including relativistic corrections to the vacuum vertex. The latter corrections yield an enhancement of order of 10 rather than the mentioned 10^2 and are therefore of minor importance compared to the larger contribution

of medium effects. The pPBF emissivity evaluated with vacuum vertex (with relativistic corrections) was implemented into their cooling code by Yakovlev et al. 2003a. Their result for this process should be still enhanced at least by a factor $\gtrsim 10$ to account for medium effects.

Medium effects essentially modify not only the pPBF emissivity but they correct also the contributions of all other processes. They have first been included into emissivities of different processes by Voskresensky & Senatorov 1984, Voskresensky & Senatorov 1986, Voskresensky & Senatorov 1987, see also Migdal et al. 1990 and a more recent review Voskresensky 2001. It was shown that the main contribution to the MU process actually comes from the *pion channel* of the reaction $nn \rightarrow npe\bar{\nu}$, where $e\bar{\nu}$ are radiated from the intermediate pion and the NN^{-1} , exchanging nucleons, rather than from the nucleon of the leg of the reaction. Moreover, due to the so called *pion softening* (medium modification of the pion propagator) the matrix elements of the MU process are further enhanced with the increase of the density towards the pion condensation critical point. Thus the corrected MU process has been called “the medium modified Urca” (MMU) and corrected NB process was called MNB. Roughly, the emissivity (1) acquires then a factor (mainly due to the pion decay channel of MMU)

$$\frac{\varepsilon_\nu[\text{MMU}]}{\varepsilon_\nu[\text{MU}]} \sim 10^3 (n/n_0)^{10/3} \frac{\Gamma^6(n)}{\omega^{*8}(n)}, \quad (5)$$

where the pre-factor $(n/n_0)^{10/3}$ arises from the phase space volume, $\Gamma(n) = 1/[1 + C(n/n_0)^{1/3}]$, $C \simeq 1.4 \div 1.6$, is the proper nucleon-nucleon correlation factor, which appears in the vertices due to the strong interaction, being expressed through the Landau-Migdal parameters (in another words $\Gamma(n)$ recovers the Ward-Takahashi identity). Note that the above introduced value Γ is an averaged quantity. Actually correlation factors depend on the energy-momentum transfer, being different for vertices connected to the weak coupling (Γ_{w-s} is rather close to unity) and for vertices related to the pure strong coupling (Γ_s is slightly less than above introduced factor Γ , $\Gamma^6 = \Gamma_{w-s}^2 \Gamma_s^4$). $\omega^*(n)$ is the so called effective pion gap. The latter has the meaning of the effective pion mass at finite momentum transfer in the given reaction, $k = p_{Fn}$. The quantity $\omega^{*2}(n)$ replaces the value $(m_\pi^2 + p_{Fn}^2)$ in the case of the free pion propagator used for the calculation of the MU process by Friman & Maxwell 1979. Thus the squared amplitude of the nucleon-nucleon interaction contains $\sim \Gamma^2/\omega^{*2}$ instead of $\sim 1/[m_\pi^2 + p_{Fn}^2]$ for the free one pion exchange used by Friman & Maxwell 1979. Here we suppressed a smaller contribution of a local Landau-Migdal interaction, being corrected by NN^{-1} loops, cf. Voskresensky & Senatorov 1986, Voskresensky & Senatorov 1987, Migdal et al. 1990.

The density dependencies of the correlation factor and the effective pion gap are presented in Fig. 1. We see that vertices are rather strongly suppressed (and this

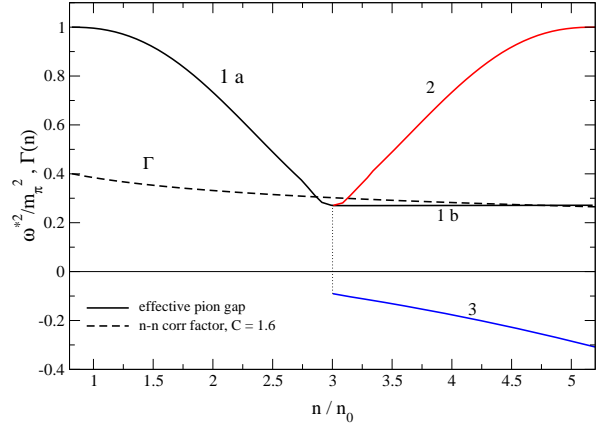


Fig. 1. Nucleon - nucleon correlation factor Γ and squared of the effective pion gap ω^* with pion condensation (branches 1a, 2, 3) and without (1a, 1b).

suppression increases with the density) but the softening of the pion mode is enhanced ($\omega^{*2} < m_\pi^2$) for $n \gtrsim 0.5 \div 0.8 n_0$. Such a behavior is motivated both theoretically and by analysis of nuclear experiments, see Migdal et al. 1990, Ericson & Weise 1988. The curve 1a shows the behavior of the pion gap for $n < n_c^{\text{PU}}$. The value n_c^{PU} depends on different rather uncertain parameters (we further assume $n_c^{\text{PU}} \simeq 3 n_0$), cf. discussion by Migdal et al. 1990. The curve 1b demonstrates the possibility of a saturation of pion softening and the absence of pion condensation for $n > n_c^{\text{PU}}$ (this possibility could be realized, e.g., if Landau-Migdal parameters increased with the density). Curves 2, 3 demonstrate the possibility of pion condensation for $n > n_c^{\text{PU}}$. The continuation of the branch 1a for $n > n_c^{\text{PU}}$ (the branch 2) demonstrates the reconstruction of the pion dispersion relation on the ground of the condensate state. Here for simplicity we do not distinguish between π^0 , π^\pm condensations, see Voskresensky & Senatorov 1984, Migdal et al. 1990 and the so called alternative-layer-structure, including both types of condensates, see Umeda et al, 1994. In agreement with a general trend known in condensed matter physics fluctuations dominate in the vicinity of the critical point of the phase transition and die out far below and above the critical point (see the curves 1a, 2). The jump from the branch 1a to 3 is due to the first order phase transition to the π condensation, see discussion of this point by Voskresensky & Mishustin 1982, Migdal et al. 1990. The branch 3 yields the amplitude of the pion condensate mean field for $n > n_c^{\text{PU}}$. The observation that the pion condensation appears by the first order phase transition needs a comment. With the first order phase transitions in the systems with several charged species is associated the possibility of the mixed phase, see Glendenning 1992. The emissivity is increased within the mixed phase since efficient DU-like processes due to nu-

cleon re-scattering on the new-phase droplets are possible. However Voskresensky et al. 2002, Maruyama et al. 2003 demonstrated that, if exists, the mixed phase is probably realized only in a narrow density interval due to the charge screening effects. Thereby to simplify the consideration we further disregard the possibility of the mixed phase. We also disregard the change in the equation of state (EoS) assuming that the phase transition is rather weak.

The works of Voskresensky & Senatorov 1984, Voskresensky & Senatorov 1986 suggested to explain the difference in the surface temperatures of various compact objects by the assumption that the objects have different masses and thus different density profiles and different cooling rates, according to above argumentation.

The most precise measurement of NS masses comes from the binary pulsar system PSRB1913+16, where for the more massive a value $M = 1.4411 \pm 0.0007 M_\odot$ has been obtained Taylor & Weisberg 1989. The approximate coincidence of this value with the present statistical average value for masses of binary pulsars $M \approx 1.35 \pm 0.04 M_\odot$ Thorsett & Chakrabarty 1999, motivated some authors to conjecture that all NS masses should be very close $1.4 M_\odot$. Recent measurement Lyne et al, 2004 of the NS mass $M \simeq 1.25 M_\odot$ in J0737-3039 does not confirm the above conjecture. NS masses can vary in some limits. Thus the argumentation, cf. Voskresensky & Senatorov 1986, that *slow to fast* cooling is explained by different masses of the corresponding objects seems rather natural. Many other in-medium reaction channels have been studied. The particular rôle of medium effects in the NS cooling has been reviewed by Migdal et al. 1990 and more recently by Voskresensky 2001. The pion softening with increasing baryon density and the subsequent pion condensation are now-days reproduced not only by microscopic models of the polarization of the baryonic medium, see Migdal et al. 1990, and Suzuki et al. 1999, but also by detailed variational calculations of Akmal et al. 1998, yielding $n_c^{\text{PU}} \sim 1.5 \div 3 n_0$ for the π^0 and charged π^- condensates. We repeat here the conclusion of the series of above mentioned works that *only due to the enhancement of medium polarization with the baryon density the pion condensate may set in and it seems thereby not justified to ignore the softening effects for $n < n_c^{\text{PU}}$, and suddenly switch on the condensate for $n > n_c^{\text{PU}}$.*

In-medium modifications of different processes including MU, NB, PFB, PU and DU reactions were implemented in the cooling code by Schaab et al. 1997 with the conclusion that the data existing to that time could be explained assuming different masses of the sources. The statement was demonstrated on different models for the equation of state (EoS) with and without superfluidity effects. However, the best fit to the whole set of the data was not elaborated. E.g., the slow cooling object PSR 1055-52 was not fitted. As we show below it cannot be done without a fit of the $1S_0$ n and p and the $3P_2$ n gaps, their absolute values and density dependencies. An alternative assumption of an internal heating suggested to explain a high temperature of this object, although possible,

seems us rather artificial and we drop it. Inclusion of different medium effects (see Voskresensky 2001) allows to develop a general “*Nuclear medium cooling scenario*” describing all the cooling data and covering the gap between “*Standard*” and “*Standard+exotics*” scenarios. The works of Blaschke et al. 2000 and Blaschke et al. 2001 considered possibility of color superconducting pure quark and hybrid stars discussing peculiarities of these objects.

The newly appeared data were studied in recent works by Kaminker et al. 2001, Yakovlev et al. 2003a and Tsuruta et al. 2002, Tsuruta 2004. Medium effects in the emissivity were disregarded. The cooling scenario based on the inclusion of the DU process (“*Standard+exotics DU*”) has been worked out by Kaminker et al. 2001, Yakovlev et al. 2003a. Their result is that the neutron $3P_2$ gap should be strongly suppressed since otherwise one could not explain the “slow cooling” objects. This assumption is supported by some modern microscopic calculations of the $3P_2$ gap including medium effects, see Schulze et al. 1996, Lombardo & Schulze 2000, Schwenk & Friman 2003. Contrary, according to Kaminker et al. 2001, Yakovlev et al. 2003a the $1S_0$ proton gap should be enhanced in order to smoothen the transition from the “slow cooling” to the “rapid cooling”. Only switching on the DU process for $M > 1.358 M_\odot$ in their EoS model allows to explain the “*intermediate cooling*” and “*rapid cooling*” of some objects. The transition from “slow cooling” to “rapid cooling” occurs in a narrow window of masses.

Criticism of the Yakovlev et al. DU “*Standard+exotics DU*” scenario is as follows, see also Tsuruta et al. 2002, Tsuruta 2004. i) The most elaborated model of the variational theory EoS of Akmal et al. 1998 permits the DU process only for rather high density $n > n_c^{\text{DU}} \simeq 5.19 n_0$ and $M \geq 1.839 M_\odot$, whereas Yakovlev et al. 2003a used a model of the EoS, where the DU process starts for $n > n_c^{\text{DU}} \simeq 2.8 n_0$ ($M > 1.358 M_\odot$). It seems doubtful that the DU reaction “might know” that it should start namely, when the NS mass is approximately $1.4 M_\odot$, being very close to the average value of the NS mass measured in NS binaries. The NS mass is governed by the full NN interaction whereas the value of the proton fraction, being responsible for the critical density of the DU reaction, is governed by only a part of the interaction related to the symmetry energy. In the light of the recent observation of the object with $M \simeq 1.25 M_\odot$ the above assumption seems even more unlikely. If the DU process occurred at $M > 1.839 M_\odot$ as follows from the EoS by Akmal et al. 1998, the model of Yakovlev et al. 2003a based on the opening of the DU channel would fail. Indeed, it would mean that the majority of the isolated NS have masses $M > 1.839 M_\odot$. The latter sounds quite unrealistic. ii) Yakovlev et al. 2003a introduced by hand a large proton gap. However, there are indications by Schulze et al. 1996 that medium effects should suppress the proton gap. Moreover, Takatsuka & Tamagaki 1997 demonstrated that proton and neutron superfluidities must be very weak for densities when the DU process be-

comes operative. iii) As we argued, the pPBF emissivity should be enhanced by an order of magnitude by medium effects, what is not taken into account. Moreover, iv) the MU and the NB processes are significantly affected by medium effects (should be replaced to MMU and MNB) what is not incorporated.

Works of Tsuruta et al. 2002, Tsuruta 2004 used the “*Standard+exotics PU*” scenario (based on the inclusion of the PU process for $n > n_c^{\text{PU}} \simeq 2.5 n_0$ and $4n_0$ respectively). In general their “*slow cooling*” curves are below those of Yakovlev et al. 2003a. Thereby the hottest object PSR 1055-52 (if the surface temperature of it is correctly extracted from the measurement) is not described, similar to the result of Schaab et al. 1997. The transition from the “*slow cooling*” to the “*rapid cooling*” is here due to the PU process. Again the model assumes that the value of the transition density for the non-standard process (here pion condensation) is very close to the value of the central density of a “magic” star of $M \simeq 1.4 M_\odot$ in order to have all NS masses very near the value of $1.4 M_\odot$. The latter seems rather unlikely in the light of the recent measurement of the NS star mass $M \simeq 1.25 M_\odot$, see Lyne et al, 2004. The medium effects are disregarded and thus the above points iii), iv) remained not corrected.

The necessity to include in-medium effects into the NS cooling problem is a rather obvious issue. It is based on the whole experience of condensed matter physics, of the physics of the atomic nucleus and it is called for by the heavy ion collision experiments, see Migdal et al. 1990, Voskresensky 2001, Ivanov et al. 2001, 1. Their relevance in the NS cooling scenario was demonstrated by Schaab et al. 1997. Below we present cooling calculations of NS based on the “*Nuclear medium cooling scenario*”, see Voskresensky 2001. We remove the above shortcomings and show the relevance of medium effects on the reaction rates and superfluid gaps by demonstrating the possibility to fit the whole set of cooling data available by today.

3. EoS and structure of NS interior, crust, surface

3.1. NS interior

We will exploit the EoS of Akmal et al. 1998 (specifically the Argonne V18 + δv + *UIX** model), which is based on the most recent models for the nucleon-nucleon interaction with the inclusion of a parameterized three-body force and relativistic boost corrections. Actually we adopt a simple analytic parameterization of this model given by Heiselberg & Hjorth-Jensen 1999 (later on HHJ). The latter uses the compressional part with the compressibility $K \simeq 240$ MeV, and a symmetry energy fitted to the data around nuclear saturation density, and smoothly incorporates causality at high densities. The density dependence of the symmetry energy is very important since it determines the value of the threshold density for the DU process. The HHJ EoS fits the symmetry energy to

the original Argonne V18 + δv + *UIX** model yielding $n_c^{\text{DU}} \simeq 5.19 n_0$ ($M_c^{\text{DU}} \simeq 1.839 M_\odot$).

Fig. 2 (left panel) shows the mass-radius relation and (right panel) the mass-central density relation. Solid lines are for HHJ EoS. In order to check an alternative possibility and to demonstrate the sensitivity of the value of the DU threshold density to the selected model we also use a version of the relativistic non-linear Walecka (NLW) model in the parameterization of Kolomeitsev & Voskresensky 2003. The parameters of the NLW model are adjusted to the following bulk parameters of the nuclear matter at saturation: $n_0 = 0.16 \text{ fm}^{-3}$, binding energy $E_{\text{bind}} = -15.8$ MeV, compression modulus $K = 250$ MeV, symmetry energy $a_{\text{sym}} = 28$ MeV, and the effective nucleon mass $m_N^*(n_0) = 0.8 m_N$. The corresponding coupling constants are as follows:

$$\begin{aligned} \frac{g_{\omega N}^2 m_N^2}{m_\omega^2} &= 91.2506 \\ \frac{g_{\sigma N}^2 m_N^2}{m_\sigma^2} &= 195.597 \\ \frac{g_{\rho N}^2 m_N^2}{m_\rho^2} &= 77.4993 \\ b &= 0.00867497, \\ c &= 0.00805981, \end{aligned} \quad (6)$$

(note that we treated obvious misprints in b and c in the paper of Kolomeitsev & Voskresensky 2003). This model was constructed in such a way that it reproduces almost the same thermodynamic properties as those in the model of Heiselberg & Hjorth-Jensen 1999. However, having not enough free parameters all NLW-based models yield significantly lower threshold densities for the DU process than those given by variational calculations. Migdal et al. 1990 treated this fact as a fragile point of the relativistic mean field models. Thereby they disregarded DU process from their subsequent analysis concentrating on the pion softening and the PU possibility. Lattimer et al. 1991 used this fact to develop the DU-based scenario for the NS cooling. Anyhow, in the given NLW model the threshold density for the DU process is $n_c^{\text{DU}} \simeq 2.62 n_0$ ($M_c^{\text{DU}} \simeq 1.29 M_\odot$). Dots in Fig. 2 (right panel) indicate threshold densities for the DU process. An influence of the pion condensation on the EoS for $n > n_c^{\text{PU}}$ is assumed to be small and suppressed thereby. For the HHJ EoS the threshold density $n_c^{\text{PU}} = 3 n_0$ corresponds to the NS mass $M_c^{\text{PU}} \simeq 1.32 M_\odot$. From Fig. 2 one can see that deviations in the $M(n)$ relation for HHJ and NLW EoS are minor, whereas the DU thresholds are quite distinct.

Fig. 3 demonstrates the concentrations of p , e and μ^- in HHJ and NLW models as a function of the baryon density. All these dependencies are quite different for these two models. This means that relativistic mean field models may describe well thermodynamic properties but yield quite different cooling picture compared with that given by more microscopically based variational calculations of Argonne-Urbana. The possibility of charged pion conden-

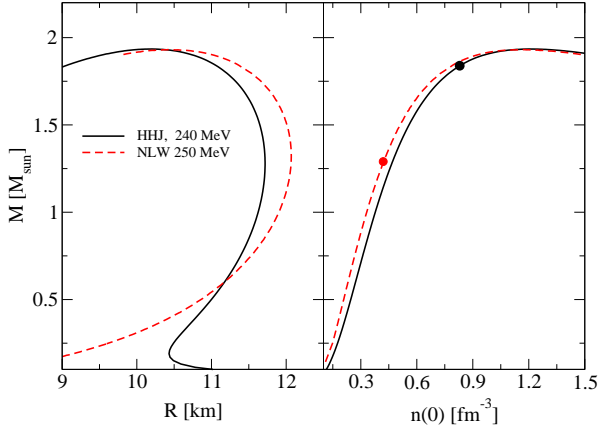


Fig. 2. Gravitational mass-radius relation and mass - central density dependence for the NS configurations corresponding to the HHJ (solid lines) and NLW (dashed lines) model EoS. Dots indicate the DU threshold. Possibility of pion condensation is suppressed.

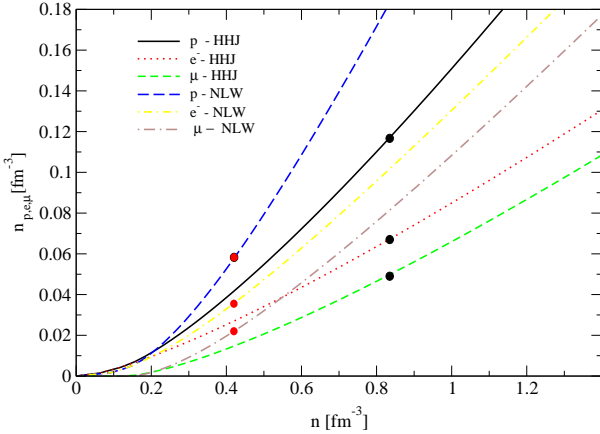


Fig. 3. Densities of the charged particles as a function of baryon density for HHJ and NLW model EoS. Threshold density for the DU process is indicated. An influence of pion condensation for $n > n_c^{\text{PU}}$ is neglected.

sation is suppressed. Otherwise for $n > n_c^{\text{PU}}$ the isotopic composition may change, see Migdal et al. 1990, in favor of an increase of a proton fraction and a smaller critical density for the DU reaction.

3.2. NS crust

The density $n \sim 0.5 \div 0.7 n_0$ is the boundary of the NS interior and the inner crust. The latter is constructed of a pasta phase discussed by Ravenhall et al. 1983, see also recent work of Maruyama et al. 2004. Then there is the outer crust and the envelope. Note that our code generates

the temperature profile being inhomogeneous during first $10^2 \div 10^3$ yr. The influence of the crust on the cooling and heat transport is rather minor basically due to its rather low mass content. Thereby the temperature also changes slightly in the crust up to the envelope.

3.3. Envelope

Further on we need the relation between the crust and the surface temperature for NS. The sharp change of the temperature occurs in the envelope. This $T_s - T_{\text{in}}$ relation has been calculated in several works, see Glen & Sutherland 1980, Yakovlev et al. 2003b, depending on the assumed value of the magnetic field at the surface and some uncertainties in our knowledge of the structure of the envelope. Fig. 4 shows a range of avail-

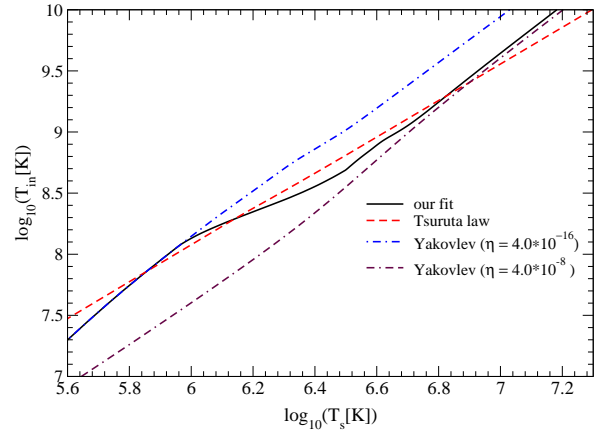


Fig. 4. Relation between the inner temperature and the surface temperature for different models used in our calculations. Dash-dotted curves indicate boundaries of possible values $T_s = f(T_{\text{in}})$, see Yakovlev et al. 2003b.

able $T_s = f(T_{\text{in}})$ curves, taken from Yakovlev et al. 2003b. The solid curve, $T_s = T_s^{\text{fit}} = f(T_{\text{in}})$, corresponds to our best fit of the “slow cooling” objects. It matches the upper boundary (the dash-dotted curve $\eta = 4.0 \cdot 10^{-16}$) for rather low T_s (in interval $\log T_s [\text{K}] \simeq 5.8 \div 6$) and the lower boundary (the dash-dotted curve $\eta = 4.0 \cdot 10^{-8}$) for high T_s ($\log T_s [\text{K}] \geq 6.8$). The dash curve shows the so called simplified “Tsuruta law” $T_s^{\text{Tsur}} = (10 T_{\text{in}})^{2/3}$ used in many old cooling calculations. We will further vary different possibilities. The dependence of the results on the $T_s - T_{\text{in}}$ relation is demonstrated below in Figs. 9 and 10, Figs. 12, 13 and 14, and Figs. 15 and 16.

4. Main cooling regulators

We compute the NS thermal evolution adopting our fully general relativistic evolutionary code. This code was originally constructed for the description of hybrid stars by

Blaschke et al. 2001. The main cooling regulators are the thermal conductivity, the heat capacity and the emissivity. In order to better compare our results with results of other groups we try to be as close as possible to their inputs for the quantities which we did not calculate ourselves. Then we add inevitable changes, improving EoS and including medium effects.

4.1. Thermal conductivity

We take the electron-electron contribution to the thermal conductivity and the electron-proton contribution for normal protons from Gnedin & Yakovlev 1995. The total contribution related to electrons is then given by

$$1/\kappa_e = 1/\kappa_{ee} + 1/\kappa_{ep}. \quad (7)$$

For $T > T_{cp}$ (normal "n" matter), we have $\kappa_{ep}^n = \kappa_{ep}$. For $T < T_{cp}$ (superfluid "s" matter), Gnedin & Yakovlev 1995 suggested to drop the superfluid contribution $1/\kappa_{ep}^s$. We use the expression

$$\kappa_{ep}^s = \kappa_{ep}/\xi_p > \kappa_{ep}^n, \quad (8)$$

that gives a crossover from the non-superfluid case to the superfluid case. The vanishing of κ_{ep}^s for $T \ll T_{cp}$ is a consequence of the scattering of superfluid protons on the electron impurities, see Blaschke et al. 2001. Following (7) we get $\kappa_e^n < \kappa_e^s$. Surprisingly, it is in disagreement with Fig. 4 of Baiko et al. 2001, where probably the curves "SF1" and "SF0" should be interchanged.

For the nucleon contribution,

$$\kappa_n = 1/\kappa_{nn} + 1/\kappa_{np}, \quad (9)$$

we use the result of Baiko et al. 2001 that includes corrections due to the superfluidity. Although some medium effects are incorporated in this work, the nucleon-hole corrections of correlation terms and the modification of the tensor force are not included. This should modify the result. However, since we did not calculate κ_n ourselves, we may only roughly estimate the modification. As we have shown above in Fig. 1, not too close to the critical point of the pion condensation the squared matrix element of the NN interaction $|M|_{\text{med}}^2 \sim p_{F,n}^2 \Gamma^2 / \tilde{\omega}^2$ is of the order of the corresponding quantity $|M|_{\text{vac}}^2 \sim p_{F,n}^2 / [m_\pi^2 + p_{F,n}^2]$ estimated with the free one pion exchange, whereas $|M|_{\text{med}}^2$ may significantly increase for $n \sim n_c^{\text{PU}}$. To simulate the effect we just allow for the variation of κ_n multiplying it by the factor $\zeta_\kappa = 10$ (fast transport) and $\zeta_\kappa = 0.3$ (slow transport), see Figs. 7, 17 below. The former case is relevant for rather massive stars, whereas the latter, for rather light stars. A suppression of the nucleon-nucleon amplitude compared to the one for the free one pion exchange model for $n \lesssim n_0$ is motivated by the in-medium T matrix calculations by Blaschke et al. 1995.

The total thermal conductivity is the straight sum of the partial contributions

$$\kappa_{\text{tot}} = \kappa_e + \kappa_n + \dots \quad (10)$$

Other contributions to this sum are smaller than those presented explicitly (κ_e and κ_n).

4.2. Heat capacity

The heat capacity contains nucleon, electron, photon, phonon, and other contributions. The main in-medium modification of the nucleon heat capacity is due to the density dependence of the effective nucleon mass. We use the same expressions as Schaab et al. 1997. The main regulators are the nucleon and the electron contributions. For the nucleons ($i = n, p$), the specific heat is (Maxwell 1979)

$$c_i \sim 10^{20} (m_i^*/m_i) (n_i/n_0)^{1/3} \zeta_{ii} T_9 \text{ erg cm}^{-3} \text{ K}^{-1}, \quad (11)$$

for the electrons it is

$$c_e \sim 6 \times 10^{19} (n_e/n_0)^{2/3} T_9 \text{ erg cm}^{-3} \text{ K}^{-1}. \quad (12)$$

Near the phase transition point the heat capacity acquires a fluctuation contribution. For the first order pion condensation phase transition this additional contribution contains no singularity, in difference with what would be for the second order phase transition, see Voskresensky & Mishustin 1982, Migdal et al. 1990. Finally, the nucleon contribution to the heat capacity may increase up to several times in the vicinity of the pion condensation point. The effect of this correction on global cooling properties is rather unimportant.

The symmetry of the $3P_2$ superfluid phase allows for the contribution of Goldstone bosons (phonons):

$$C_G \simeq 6 \cdot 10^{14} T_9^3 \frac{\text{erg}}{\text{cm}^3 \text{ K}}, \quad (13)$$

for $T < T_{cn}(3P_2)$, $n > n_{cn}(3P_2)$. We also include this contribution in our study although its effect on the cooling is rather minor.

4.3. Emissivity

We adopt the same set of partial emissivities as in the work of Schaab et al. 1997. The phonon contribution to the emissivity of the $3P_2$ superfluid phase is negligible. The main emissivity regulators are the MMU, see above rough estimation (5), nPBF and pPBF processes, see above rough estimation (4).

Only qualitative behavior of the interaction shown in Fig. 1 is motivated by microscopic analysis whereas actual numerical values of the correlation parameter and the pion gap are rather uncertain. Thereby we vary the values $\Gamma(n)$ and $\omega^{*2}(n)$ in accordance with our above discussion of Fig. 1. By that we check the relevance of alternative possibilities: a) no pion condensation and a saturation of the pion softening with increasing density, b) presence of pion condensation.

We also add the contribution of the DU for $n > n_c^{\text{DU}}$, see above rough estimation (3).

All emissivities are corrected by correlation effects. The PU process contains an extra Γ_s^2 factor compared to the DU process. Another suppression of PU emissivity comes from the fact that it is proportional to the squared pion condensate mean field $|\varphi|^2$. Near the critical point $|\varphi|^2 \sim 0.1$ increasing with the density up to $|\varphi|^2 \sim f_\pi^2/2$,

where $f_\pi \simeq 93$ MeV is the pion decay constant. Finally, the PU emissivity is about 1-2 orders of magnitude suppressed compared to the DU one. Moreover, we adopt the same gap dependence for the PU process as that for the DU process.

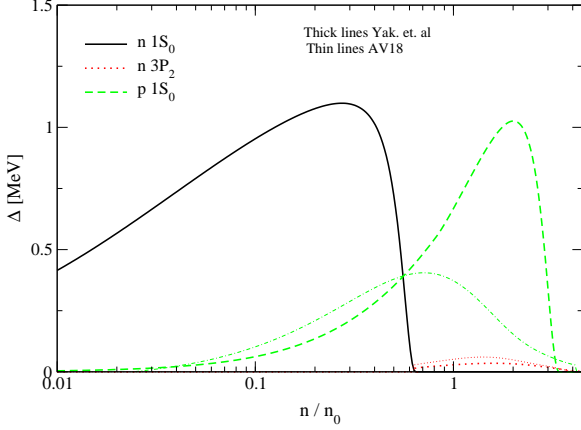


Fig. 5. Neutron and proton pairing gaps according to Yakovlev et al. 2003a (thick solid, dashed and dotted lines) and to Takatsuka & Tamagaki 2004 (thin lines).

4.4. Nucleon superfluidity

In spite of many calculations, values of nucleon gaps in dense NS matter are poorly known. This is the consequence of the exponential dependence of gaps on the poorly known NN interaction for $n \neq n_0$. Recent calculations of Schwenk & Friman 2003 who included medium effects in evaluation of the $3P_2$ gap demonstrate its strong suppression up to values $\lesssim 10$ keV. Together with previous findings of Schulze et al. 1996, Lombardo & Schulze 2000 they motivate to consider possibility of rather suppressed gaps. The suppression of the $3P_2$ gap for $n > n_c^{\text{PU}}$ (for charged π condensation), as explained in text of the work of Schaab et al. 1997, is also in coincidence with the argumentation of Takatsuka & Tamagaki 1997, see also Tsuruta et al. 2002, Tsuruta 2004.

Below we start with the model used by Yakovlev et al. 2003a. Also we use the gaps from recent work Takatsuka & Tamagaki 2004 to compare the results. The gaps are presented in Fig. 5. Thick lines, the gaps from Yakovlev et al. 2003a and thin lines, from Takatsuka & Tamagaki 2004, the model AV18 by Wiringa et al. 1995. The $1S_0$ neutron gap is taken the same in both models. We allow for a variation of gaps in wide limits in order to check the sensitivity of the results to their values. As we will see the cooling curves are sensitive to the magnitudes and to the density dependence of the gaps. Therefore further microscopic studies of the gaps are required.

4.5. KDU, HDU, QDU and other

The phase structure of dense NS matter might be very rich, including π^0 , π^\pm condensates and \bar{K}^0 , K^- condensates in both S and P waves (Kolomeitsev & Voskresensky 2003); charged ρ -meson condensation by (Voskresensky 1997); coupling of condensates Umeda et al, 1994; fermion condensation yielding an efficient DU-like process in the vicinity of the pion condensation point (with the emissivity $\varepsilon_\nu \sim 10^{27} T_9^5$, $m_N^* \propto 1/T$), see Voskresensky et al. 2000; hyperonization, see Takatsuka & Tamagaki 2004; quark matter with different phases, like so called 2SC, CFL, CSL, plus their interaction with meson condensates, see K. Rajagopal & F. Wilczek 2000, Blaschke et al. 2001 and refs therein; and different mixed phases. In the present work, we suppress all these possibilities of extra efficient cooling channels. They are effectively simulated by our PU choice. The quark matter effects need a special discussion. We will return to the latter possibility in a subsequent publication.

5. Numerical results

5.1. Cooling of normal NS

Although we have no doubts about the presence of nucleon superfluidity in NS interiors, we consider first the case of the complete absence of nucleon superfluidity. The reasons for that are as follows: i) Thus we compare our results with previous calculations. ii) On this example we select more essential and less essential ingredients in this many-parameter problem. Results are more transparent being demonstrated on a simplified example since the general case introduces new uncertain parameters and is more involved. iii) The actual values of gaps might be essentially smaller than those estimated in the literature, cf. Lombardo & Schulze 2000, Schwenk & Friman 2003. This is because most calculations did not include a proper medium dependence of the NN interactions. Thus we discuss the limiting case of largely suppressed gaps.

Fig. 6 demonstrates the cooling evolution (for $T_s^{\text{fit}}(t)$ dependence) of a normal NS for the HHJ EoS. Medium effects and PU possibility are disregarded (“*Standard scenario*”). We see that one could easily explain the “*slow cooling*” points. Curves for NS masses in the range $M \simeq (1 \div 1.839) M_\odot$ lie in this region. For $M \simeq 1.839 M_\odot$ the efficient DU process is switched on within the HHJ EoS (*Standard + exotics DU* scenario). Thereby, the curve corresponding to $M \simeq 1.841 M_\odot$ jumps down and already explains the “*rapid cooling*” points. The “*intermediate cooling*” points can be explained either by a very low NS mass (see curve for $M \simeq 0.5 M_\odot$, or by $M \simeq 1.840 M_\odot$. However, it seems rather unrealistic that 4 from 10 objects either relate to very low masses like $M \simeq 0.5 M_\odot$, or are highly massive, as $M \simeq 1.840 M_\odot$, where the DU process is switched on. “*Rapid cooling*” points are explained by very massive objects ($M \simeq 1.841 M_\odot$) with DU process. They could be also explained by very low-mass objects

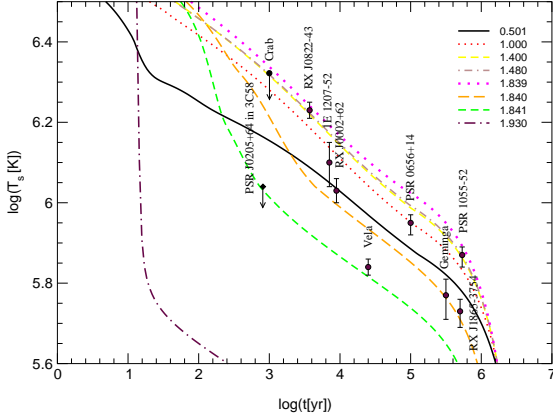


Fig. 6. Cooling of HHJ configurations without medium effects and PU for different masses of NS, $T_s - T_{\text{in}}$ relation according to our fit.

($M \sim 0.1 M_\odot$). We further drop the latter possibility as rather unrealistic one and consider $M \geq 0.5 M_\odot$.

Thus the picture as a whole looks unsatisfactory. *Only “slow cooling” data are appropriately explained in a reasonable NS mass interval $1.0 \div 1.839 M_\odot$.* The explanation of “intermediate cooling” and “rapid cooling”, although possible, needs very unnatural assumptions.

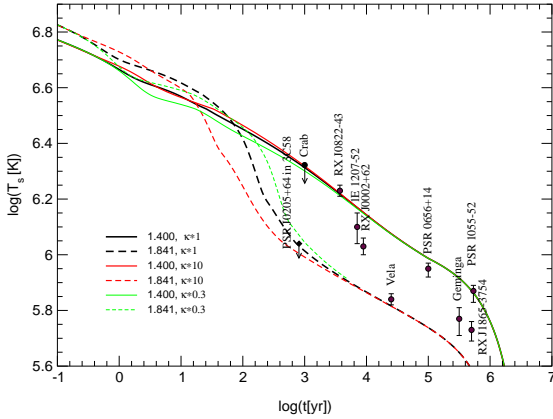


Fig. 7. The influence of the heat conductivity on the scenario of Fig. 6. Two representative configurations with masses $M = 1.40 M_\odot$ (solid lines) and $M = 1.841 M_\odot$ (dashed lines). We vary the heat conductivity by factors $\zeta = 10$ and $\zeta = 0.3$.

Fig 7 demonstrates the sensitivity of the cooling curves to the variation of κ_n . We scale κ_n by a factor $\zeta = 10$ and by $\zeta = 0.3$. The former case is meaningful for heavy objects whereas the latter one for low-mass objects. We see that *both increasing and decreasing of the thermal conduc-*

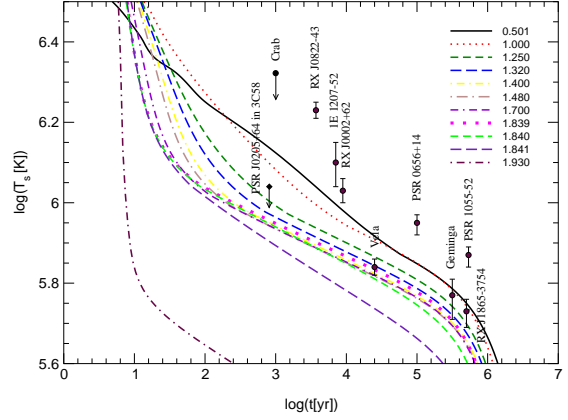


Fig. 8. Cooling evolution of the NS with normal matter HHJ EoS (for T_s^{fit}) including the medium modification of MU and other processes (MMU, MNB, etc.), without pion condensation.

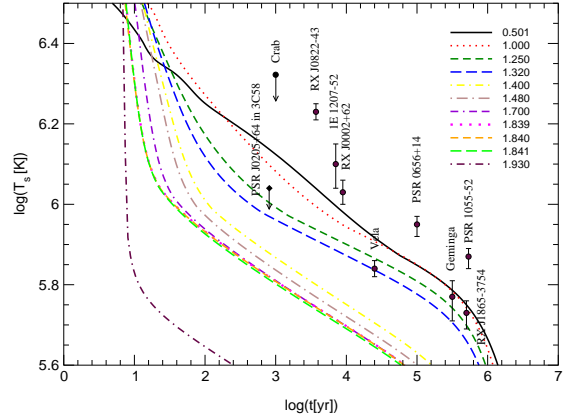


Fig. 9. Same as Fig. 8, including pion condensation for $n > 3 n_0$.

tivity does not change the picture as the whole, as well as the conclusion drawn above. Transport is relevant only up to the first 10^3 y when the details of the $T_s - T_{\text{in}}$ relation as well as temperature and density dependences of the cooling regulators determine the response of the cooling curves to a rescaling of the heat conductivity by the factor ζ .

Fig. 8 allows for medium effects the strength of which is assumed to be saturated with increasing density (no π condensation for $n > 3 n_0$). Thus we use curves 1a, 1b of Fig 1. We see that medium effects being included in calculation of the emissivity (MMU, MNB and others) significantly decrease all the curves. This allows us to explain reasonably well “rapid cooling” points but “slow cooling” and “intermediate cooling” cannot be addressed without inclusion of superfluidity. A partial suppression of

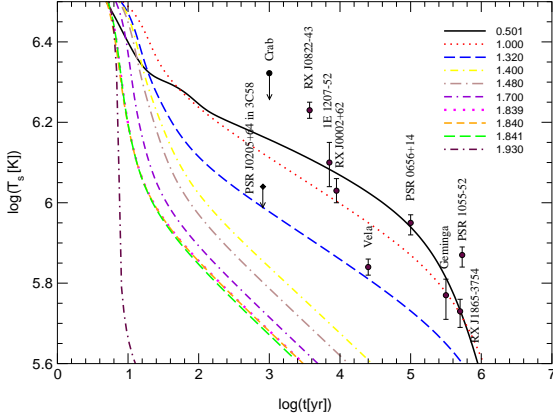


Fig. 10. Same as Fig. 9, but with $T_s - T_{in}$ relation given by the Tsuruta law.

medium effects by scaling of ω^* does not allow to improve the picture.

Comparison of Fig. 8 and Fig. 6 shows that inclusion of medium effects regulates the mass dependence of the curves. In Fig. 6 the curves rise with increase of the NS mass (for objects with $M \leq 1.89 M_\odot$, below DU threshold) whereas in Fig. 8 the trend is changed to opposite.

Fig. 9 additionally allows for π condensation for $n > 3n_0$ ($M > 1.32 M_\odot$ for HHJ EoS). Now we use curves 1a, 2, 3 of Fig. 1. The picture remains almost the same as that shown in Fig. 8. However objects with $M \geq 1.32 M_\odot$ are cooled still faster by the efficient PU reaction. If we assumed $n_c^{PU} \simeq 2.5 n_0$ the pion condensation would start for NS masses $M > 1.08 M_\odot$.

Fig. 10 demonstrates the same cooling evolution of normal NS as Fig. 9, but for the Tsuruta relation $T_s^{Tsur}(t)$. One can explain the “rapid cooling” and the “intermediate cooling” but cannot explain the “slow cooling”. We checked that the use of other $T_s - T_{in}$ relations in the frame given by Fig. 4 does not change general trends. *Medium effects and the possibility of π condensation, although they regulate the NS mass-dependence show too rapid cooling.* We see that without superfluidity the picture is unsatisfactory. Thereby we conclude that *the cooling data call for the nucleon superfluidity.*

5.2. Cooling of superfluid NS

First we check the best fit of the model Kaminker et al. 2001, Yakovlev et al. 2003a, where gaps are given by Fig. 5 (thick lines). In order to get a fit of the data within their “Standard + DU” scenario Kaminker et al. 2001, Yakovlev et al. 2003a were forced to additionally switch off the $3P_2$ neutron gap. Since the full switching off the $3P_2$ gap is not supported by microscopic calculations we simulate the same effect by introducing the scaling factor 0.1 for the $3P_2$ neutron

gap shown in Fig. 5 by corresponding solid line. Then the magnitude of the $3P_2$ gap becomes to be comparable with the value following from the findings of Schwenk & Friman 2003, who estimated the $3P_2$ neutron gap including medium effects.

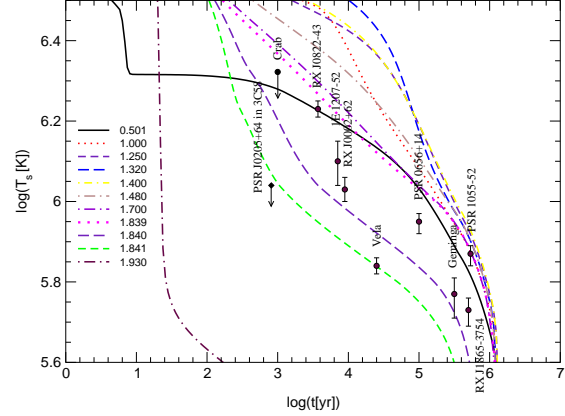


Fig. 11. Cooling of NS configurations with superfluid nuclear matter without medium effects and pion condensation. The gaps are taken as in Yakovlev et al. 2003a, see Fig. 5, the neutron $3P_2$ gap is additionally suppressed by a factor 10, $T_s - T_{in}$ relation is given by $T_s^{\text{fit}}(t)$.

Fig. 11 demonstrates the cooling evolution ($T_s^{\text{fit}}(t)$ dependence) of the above choice of gaps for a NS with HHJ EoS. Medium effects are not included. The curves rise compared to those in Fig. 6. We see that for $M \leq 1.839 M_\odot$ when the DU process is switched off, all the curves demonstrate very slow cooling even not explaining the “slow cooling” points (for $M \gtrsim 1 M_\odot$). We also checked the artificially suppressed rate of the pPBF process used by Kaminker et al. 2001, Yakovlev et al. 2003a. Suppression of the rate by a factor of 10 does not significantly affect the curves and does not change the conclusion. The general trends of the curves are similar to those of Kaminker et al. 2001. However within the HHJ EoS the “intermediate cooling” points and “rapid cooling” points can be explained only by objects with $M > 1.839 M_\odot$ and DU cooling what seems us quite unsatisfactory. Comparison of Fig. 11 with Fig. 6, computed within the same scenario but without inclusion of the superfluidity, shows that both choices suffer of the very same shortcomings. *The data call for medium effects.*

Fig. 12 allows for the medium effects (we use curves 1a, 1b of Fig. 1, no π condensation for $n > 3n_0$). We smoothly cover all data points. All the data are explained by masses in the interval $M \simeq 1.39 \div 1.84 M_\odot$. The regular behavior is clearly seen. More massive objects cool faster than less massive ones. The star with the mass $1.25 M_\odot$ may relate to the old and hot objects like PSR 1055-52.

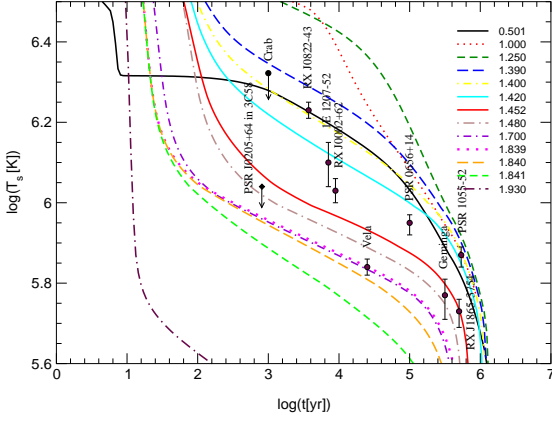


Fig. 12. Same as in Fig. 11 including medium effects without pion condensation (curves 1a, 1b in Fig. 1).

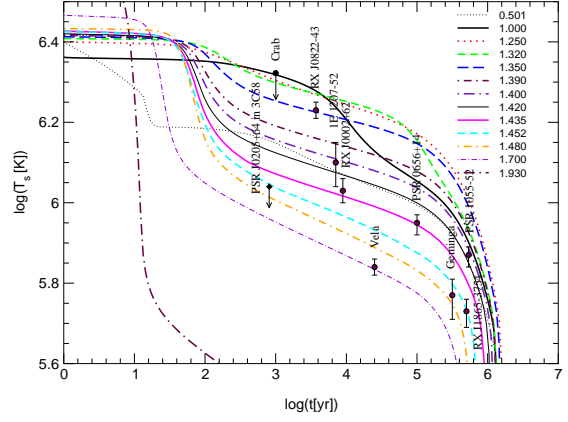


Fig. 14. Same as in Fig. 12 but $T_s - T_{in}$ relation according to Yakovlev et al. 2003a, $\eta = 4 \times 10^{-16}$.

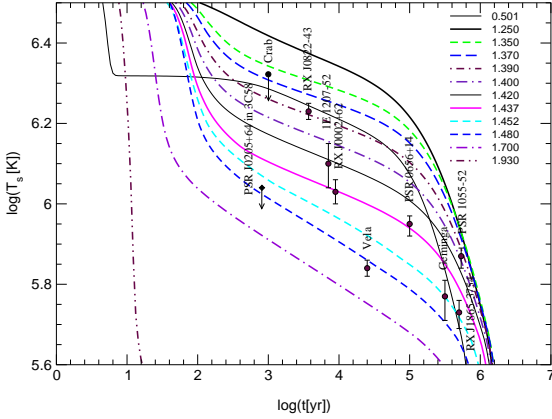


Fig. 13. Same as in Fig. 12 but using the Tsuruta law.

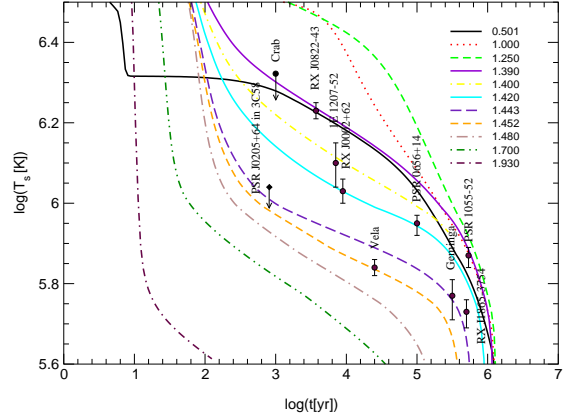


Fig. 15. Same as in Fig. 12 including π condensation.

We checked a sensitivity of the result to the dependence $T_s - T_{in}$. Fig. 13 uses the “Tsuruta law” whereas Fig. 14 probes the choice of Yakovlev et al. 2003a, $\eta = 4.0 \times 10^{-16}$. We see that a variation of $T_s - T_{in}$ does not change the picture as the whole. Only the interval of masses that cover the data is slightly changed. This interval is $M \simeq 1.36 \div 1.50 M_\odot$ according to Fig. 13 and $M \simeq 1 \div 1.75 M_\odot$ according to Fig. 14.

The mass interval that covers the data is $M \simeq 1.37 \div 1.46 M_\odot$. Fig. 15 allows for π condensation for $n > 3 n_0$. Now we use curves 1a, 2, 3 of Fig 1. The $T_s - T_{in}$ relation is according to our fit.

Fig. 16 shows the same as Fig. 15 but for the Tsuruta law. In both cases the picture remains almost the same as that shown in Fig. 13 and Fig. 12, respectively. The difference starts for masses $\geq 1.32 M_\odot$ due to switching on the efficient PU process. The interval of masses that cover the data is $M \simeq 1.36 \div 1.452 M_\odot$. Using of other

$T_s - T_{in}$ dependences does not change the picture. Thus we may conclude that the π condensation does not contradict the data but, on the other hand, this assumption is not motivated by the existing cooling data. The hyperon enhanced cooling or kaon condensation cooling do not change the picture if their critical densities are $\gtrsim 3 n_0$. Note that the transition density can’t be too low. At least $n_c^{PU} \geq 2.5 \div 2.7 n_0$ for the charged pion condensation within given EoS. The neutral PU process might be additionally suppressed compared to the charged one, see Voskresensky 2001, that may weaken the above restriction. Otherwise we would get too fast cooling already for low mass objects and the NS with $M \simeq 1.4 M_\odot$ would cool faster than it is available by modern data.

Fig. 17 demonstrates the sensitivity of the curves presented in Fig. 15 to the variation of κ_n . We scale κ_n by the factor $\zeta = 10$ and $\zeta = 0.3$. We see that both increasing and decreasing of the thermal conductivity κ_n does not

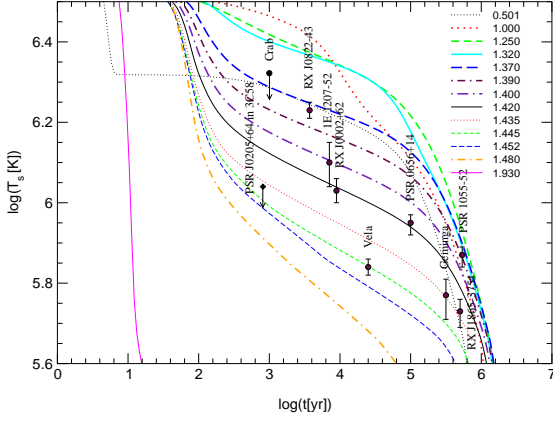


Fig. 16. Same as in Fig. 15 using the Tsuruta law.

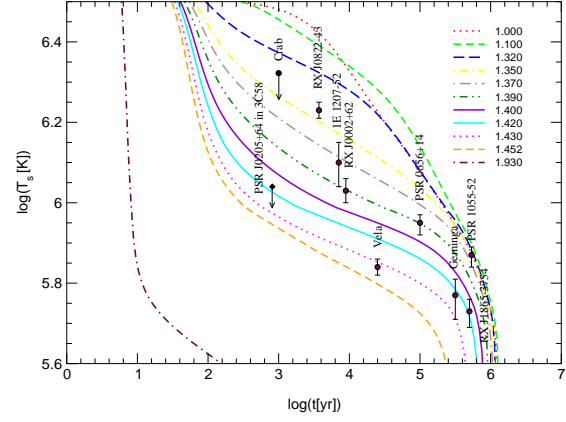


Fig. 18. Same as Fig. 15 but with suppressed proton gap by a factor 0.5.

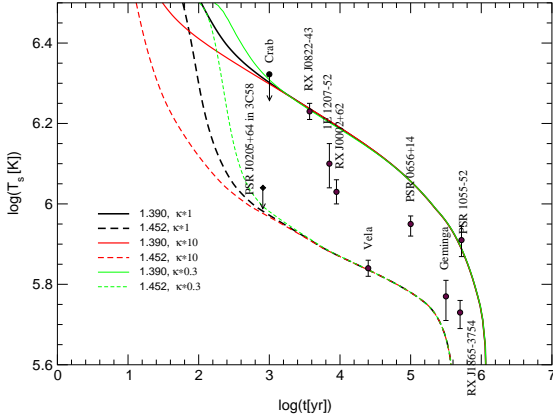


Fig. 17. The influence of a change of the heat conductivity on the scenario of Fig. 15.

change the picture as the whole as well as the conclusion drawn above. As in the case when medium effects are suppressed, see Fig. 7, the effect of the transport is relevant for first 10^3 yr.

Figs. 18 - 23 show the dependence of the results on different variations of the gaps, their absolute values and the density dependencies.

We remind the reader that the proton gap is artificially enhanced by Yakovlev et al. 2003a. Fig. 18 presents the same as Fig. 15, but now we suppress the proton gap by a factor 0.5. An appropriate fit is achieved. In Fig. 19 we further suppress the proton gap (now by factor 0.2) and we also suppress the $1S_0$ neutron gap by factor 0.5 simulating the medium effects in gaps. We again obtain an appropriate overall fit of the data. The NS masses covering available data are $1.33 \div 1.44 M_\odot$ in case of Fig. 18 and $1.23 \div 1.42 M_\odot$ in case of Fig. 19. However we have checked that scaling of all gaps shown by thick lines in Fig. 5 by

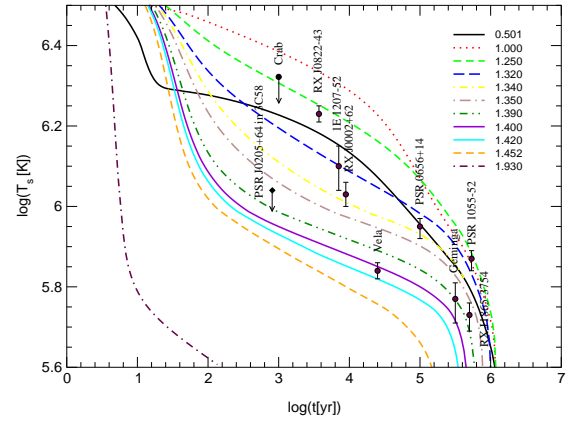


Fig. 19. Same as Fig. 15 but with $1S_0$ neutron gap suppressed by a factor 0.5 and $1S_0$ proton gap by a factor 0.2. The neutron $3P_2$ gap is remained to be suppressed by 0.1.

the factor 0.1 already does not allow for the appropriate fit of the data. Thus we demonstrated important role played by all three types of the pairing: $1S_0$ nn and pp and $3P_2$ nn .

In Fig. 20 we include the same medium effects as those in Fig. 12 but we use the gaps according to Takatsuka & Tamagaki 2004, see thin lines in Fig. 5, additionally suppressing $3P_2$ gap by a factor 10. We see that with HHJ EoS and the given choice of the gaps we obtain *the best fit of the whole set of the data*. Thus the pion softening and appropriate density dependence of the gaps are quite sufficient to explain the modern NS cooling data. The data are covered by NS in wide mass interval $0.5 \div 1.84 M_\odot$. For masses $1.0 \div 1.84 M_\odot$ the picture is quite regular. Less massive stars cool slower, more massive stars cool faster.

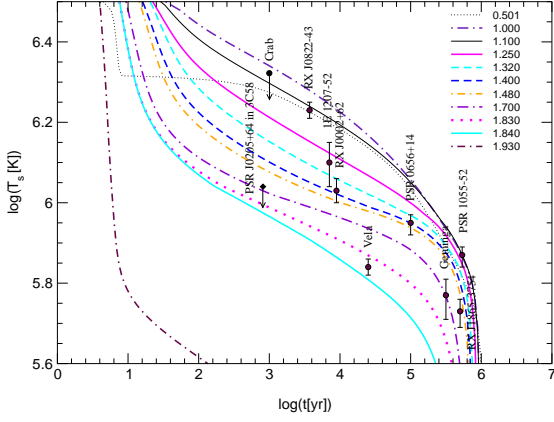


Fig. 20. Same as in Fig. 12 using the gaps of Takatsuka & Tamagaki 2004, see Fig. 5, with additional suppression of the $3P_2$ gap by a factor 10.

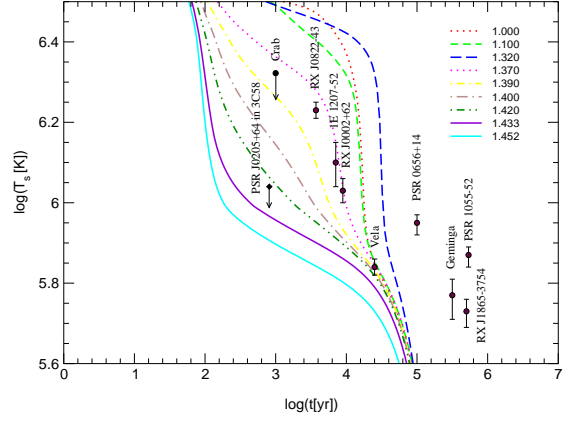


Fig. 22. Same as Fig. 15 without suppression of $3P_2$ neutron pairing gap.

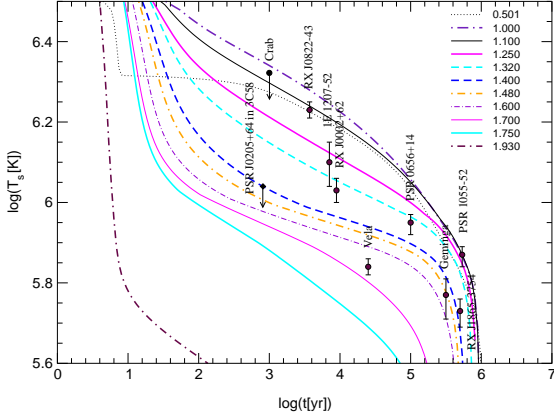


Fig. 21. Same as in Fig. 20 but including pion condensation.

Fig. 21 presents the same as Fig. 20 but now including the pion condensation. The picture as the whole remains the same. Only cooling of massive stars is touched that slightly narrows the appropriate NS mass interval from $1.0 \div 1.75 M_\odot$ to $1.0 \div 1.70 M_\odot$.

Fig. 22 and Fig. 23 present the same as Fig. 15 and Fig. 21, but now we take off the suppression factor 0.1 for the $3P_2$ neutron gap. From Fig. 22 we see that 2 ‘slow cooling’ and 2 ‘intermediate cooling’ points related to old NS are not explained. In case of Fig. 23 only two lowest data points are explained. We again support the statement that *the suppression of the $3P_2$ gap is indeed called for by the cooling data.*

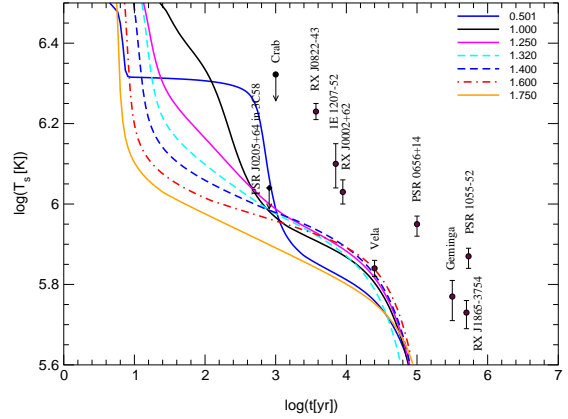


Fig. 23. Same as Fig. 21 without suppression of $3P_2$ neutron pairing gap.

6. Conclusion

We have shown that the most up-to-date observed NS cooling data can well be explained. Actually we deal with a many-parametric problem that allows for a variation of many quantities. Besides the data points might be partially shifted from their positions shown in figures due to a number of uncertainties and assumptions used at their analysis. By above figures we have illustrated different possibilities discriminating more probable explanations from less probable ones. We elaborated the example of the probably most realistic and microscopically supported EoS of the $V18 + \delta v + UIX^*$. Actually we used a simple parameterization of this EoS suggested by Heiselberg & Hjorth-Jensen 1999 (HHJ). In this EoS the DU process does not show up to $M \simeq 1.839 M_\odot$. Thus within this model the ‘Standard + DU’ scenario would demonstrate that the majority of experimentally

measured cooling points relate to very massive NS, with $M \geq 1.84 M_{\odot}$. From our point of view such a scenario seems unrealistic.

We exploit in-medium effects in the calculation of the emissivities and the pairing gaps, and, that is less important, in specific heat and heat conductivity. In general, medium effects result in a significant suppression of the superfluid gaps, especially of the $3P_2$ nn gap, and they enhance the cooling rates of the MMU and the MNB processes through *the pion softening effect with the increase of the density*. Without the latter effect the “Standard + PU” scenario suffers from an internal inconsistency. Pion condensation cannot take place without preliminary softening of the pion mode at lower densities. And v.s., recent argumentation for the pion condensation, cf. Akmal et al. 1998, Suzuki et al. 1999, further motivates the presence of a precursor pion softening, details see in Migdal et al. 1990. Pion condensation at $n \gtrsim 3n_0$ does not contradict the data but the data can be explained also without pion condensation and other so called “exotics” (KU, HDU, etc) but with the pion softening. Based on the works Schaab et al. 1997, Blaschke et al. 2001 we have further improved our code. We showed that the results are sensitive to the values and the density dependencies of the nucleon superfluid gaps. *The strong suppression of the $3P_2$ gap* motivated by theoretical evaluations that incorporate medium effects is indeed required for the fit of the data. We show that all three groups of points “slow cooling”, “intermediate cooling” and “rapid cooling” are now well explained on the basis of the “Nuclear medium cooling scenario” demonstrated here, where an important rôle is played by in-medium effects, cf. Voskresensky 2001.

We may draw the following main conclusions.

i) The normal matter assumption (see Figs. 6 - 10) seems rather unrealistic, as by itself, as in relation to the cooling data. One could explain the data but at the price of ignoring of medium effects in MMU and MNB. Then the “intermediate cooling” points can only be explained by very low NS masses (as $0.5 M_{\odot}$) or, together with “rapid cooling” points, by the very high NS masses $M > 1.839 M_{\odot}$, which allow for the DU process. In the latter case the mass window separating the “intermediate cooling” and “rapid cooling” is very narrow. Above price seems us too high and we drop such a scenario. The superfluidity is called for by the data.

ii) Including superfluid gaps we see, in agreement with recent microscopic findings, that $3P_2$ neutron gap should be as small as $\lesssim 10$ keV. Otherwise one cannot explain at least several old objects (see Figs. 22, 23). Proton and $1S_0$ neutron gaps also might be suppressed by factors of order of several, as it is demonstrated by Figs. 18, 19, 20, 21, but the suppression by factor $\gtrsim 10$ is not already permitted.

iii) Medium effects associated with the pion softening are called for by the data. As the result of the pion softening the pion condensation may occur for $n \geq n_c^{\text{PU}}$ ($n \geq 3n_0$ in our model). Its appearance does not contradict to the data (see Fig. 15) but also the data are well described, if the softening effect is rather saturated (see

Fig. 12) with increase of the density (as demonstrated by curves 1a, 1b in Fig. 1). At the same time the critical densities for the efficient DU-like exotic processes (such as PU on charged pion) should not be as small as $< 2.5 n_0$, if the given HHJ EoS is indeed correct. Otherwise one would get too rapid cooling of the NS of the typical $1.4 M_{\odot}$ mass. This also means that the proper DU threshold density can’t be too low that puts restrictions on the density dependence of the symmetry energy. Both statements might be important in the discussion of the heavy ion collision experiments.

iv) We demonstrated a regular mass dependence: for the NS masses $M \gtrsim 1.0 M_{\odot}$ less massive NS cool slower, more massive NS cool faster.

As we have mentioned, for the sake of simplicity we did not include the possibility of the hyperonization and other possibilities, like kaon condensation and fermion condensation, which may stimulate a more rapid cooling, being working in a line with the pion condensation. We did not include possible quark effects. The latter need a special treatment. The possibility of the color superconductivity in dense NS interiors opens a number interesting possibilities like the so called two-flavor color superconductivity (2SC) phase, color-flavor-locking (CFL) phase, color-spin-locking (CSL) phase. Their cooling is essentially different. We will return to this discussion in the nearest future.

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